

Light Higgs from Scalar See-Saw in Technicolor

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We consider a TeV scale see-saw mechanism leading to light scalar resonances in models with otherwise intrinsically heavy scalars. The mechanism can provide a 125 GeV technicolor Higgs in *e.g.* two-scale TC models.

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I. INTRODUCTION

Recently the ATLAS and CMS collaborations have announced the discovery of a new boson with a mass of approximately 125 GeV [1, 2]. The next challenge is to determine the nature of this new state, including its quantum numbers and couplings, and whether it is fundamental or composite. Because of the observed decays to $\gamma\gamma$, WW , and ZZ , with strengths seemingly comparable to those expected from the standard model (SM) Higgs, this state is likely a (mostly) CP -even spin-zero boson, although CP -odd candidates are not currently ruled out [3–7].

In technicolor (TC), *common lore* has it that the lightest CP -even spin-zero resonance, the analogue of the σ meson or $f_0(500)$ in QCD, cannot be as light as 125 GeV. In fact TC theories are sometimes considered as underlying theories for Higgsless models, despite the fact that in QCD the σ meson is among the lightest states. In this paper we consider the possibility of a light TC Higgs arising from mass mixing between relatively heavy scalar singlets. This results in a “see-saw” mechanism, with one scalar singlet becoming lighter and one heavier than the corresponding diagonal mass. This is expected to occur in two-scale TC models, *e.g.* low-scale TC [8, 9] and ultra-minimal TC (UMT) [10]¹.

Two-scale TC theories feature two technifermion species with different representations under a single technicolor gauge group. These lead, for instance, to two different sets of composite scalars. Because of different quantum numbers, scalar multiplets from different representations do not mix through mass terms. However scalar singlets do. Because of the strength of the TC interaction, such a mixing can be sizable, which is the key ingredient in the see-saw mechanism. Moreover, radiative corrections from the top quark may contribute to further reduce the mass of the lightest scalar singlet [13].

This paper is organized as follows. In Sec. II we briefly review the see-saw mechanism for scalar singlets. In Sec. III we review the spin-zero sector of two-scale TC, and analyze the properties of the mixing mass term. Then we apply the general results to UMT and low-scale TC. Finally, in Sec. IV we offer a brief discussion of our findings.

¹ In models with fundamental scalars, a scalar see-saw mechanism has also been considered as a way of generating a negative mass squared for the Higgs [11, 12].

II. HIGGS SEE-SAW MECHANISM

Consider a theory featuring two scalar singlets in its spectrum, H_1 and H_2 . Assume these to mix via mass term:

$$\mathcal{L} \supset -\frac{M_1^2}{2}H_1^2 - \frac{M_2^2}{2}H_2^2 - \delta M_1 M_2 H_1 H_2. \quad (1)$$

In the limit $\delta^2 \rightarrow 1$ one eigenstate is massless. It is therefore useful to define the parameter

$$\varepsilon \equiv 1 - \delta^2. \quad (2)$$

Diagonalization gives

$$\begin{pmatrix} H_1 \\ H_2 \end{pmatrix} = \begin{pmatrix} \cos \beta & \sin \beta \\ -\sin \beta & \cos \beta \end{pmatrix} \begin{pmatrix} H_- \\ H_+ \end{pmatrix}, \quad \tan 2\beta = \frac{2M_1 M_2}{M_2^2 - M_1^2} \delta, \quad (3)$$

where H_- and H_+ are the light and heavy mass eigenstate, respectively, with mass

$$M_{\pm}^2 = \frac{M_1^2 + M_2^2}{2} \left[1 \pm \sqrt{1 - \left(\frac{2M_1 M_2}{M_1^2 + M_2^2} \right)^2 \varepsilon} \right]. \quad (4)$$

For $\varepsilon \ll 1$, M_-^2 becomes

$$M_-^2 = \frac{M_1^2 M_2^2}{M_1^2 + M_2^2} \varepsilon + \mathcal{O}(\varepsilon^2) \ll M_1^2, M_2^2. \quad (5)$$

In Fig. 1 we show M_- (solid) and M_+ (dashed), for the cases $M_1 = M_2 = 1.0$ TeV (black) and $M_1 = 1.0$ TeV, $M_2 = 300$ GeV (red), as a function of $|\delta|$. The dotted horizontal line corresponds to the experimental value of 125 GeV. This trivial exercise illustrates how theories with relatively heavy scalar mass scales, *e.g.* M_i between a few hundreds of GeVs and a few TeVs, may still feature a light scalar eigenstate after diagonalization, and thus a candidate for the recently observed 125 GeV boson. While this mechanism is simple and general, it is immediately applicable to TC models with two condensation scales. In the next section we discuss the general properties of singlet scalar mixing in two-scale TC theories, and provide specific examples.

III. TWO-SCALE TECHNICOLOR

In two-scale TC two dynamical scales arise due to the presence of two Dirac technifermion species Q_i , transforming under different representations R_i of a single TC gauge group [8, 14]. The TC force causes technifermion bilinears to condense at different scales Λ_i . An estimate of Λ_2/Λ_1

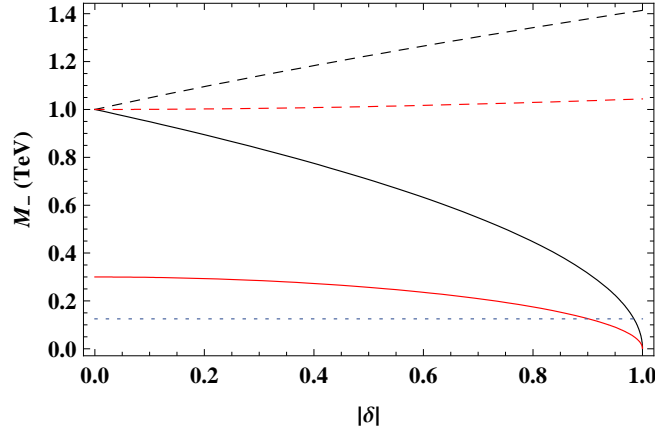


FIG. 1: M_- (solid) and M_+ (dashed) as a function of $|\delta|$, for the cases $M_1 = M_2 = 1.0$ TeV (black) and $M_1 = 1.0$ TeV, $M_2 = 300$ GeV (red). The dotted horizontal line corresponds to the experimental value of 125 GeV.

can be obtained from the ladder Schwinger-Dyson equation for the techniquark propagator. In this approximation, the critical coupling for chiral symmetry breaking depends on the representation via $\alpha_c(R_i) = \pi/3C_2(R_i)$, where $C_2(R_i)$ is the Casimir of the representation R_i . Taking $C_2(R_1) \leq C_2(R_2)$, and integrating the one-loop beta-function $\beta(\alpha) = -\beta_0(R)\alpha^2/2\pi$ from Λ_1 to Λ_2 gives

$$\frac{\Lambda_2}{\Lambda_1} \simeq \exp \left[\frac{2\pi}{\beta_0(R_1)} \left(\alpha_c(R_2)^{-1} - \alpha_c(R_1)^{-1} \right) \right]. \quad (6)$$

Since $\Lambda_1 \leq \Lambda_2$, or equivalently $\alpha_c(R_1) \geq \alpha_c(R_2)$, the fermions in the representation R_2 are effectively decoupled below Λ_2 . Therefore, only $\beta_0(R_1)$ appears in the exponent. If $\beta_0(R_1)$ and $\alpha_c(R_1)$ are small then the scale separation can be sizeable and the presence of four-fermion operators can contribute to further enhance the scale separation [15]. This crude approximation serves to illustrate the appearance of two distinct scales.

Now, let N_i be the number of Dirac techniflavors in the representation R_i . The global symmetries of the corresponding fermion sector depend on whether R_i is complex, real, or pseudoreal. To see this we express the Dirac fermions in terms of two Weyl fermions,

$$Q_{im} = \begin{pmatrix} \psi_{im}^1 \\ \bar{\psi}_{im}^2 \end{pmatrix}, \quad (7)$$

where m is the techniflavor index ($m = 1, \dots, N_i$), and the color index is suppressed. If ψ_{im}^1 transforms under the R_i representation, then $\bar{\psi}_{im}^2$ transforms under the conjugate representation \bar{R}_i . For complex R_i this implies that rotations in techniflavor space *cannot* mix ψ^1 and ψ^2 fermions.

As a consequence the TC Lagrangian features a global $SU(N_i)_1 \times SU(N_i)_2 \times U(1)$ techniflavor symmetry, where the extra $U(1)$ corresponds to technibaryon-number conservation. This global symmetry is spontaneously broken to diagonal $SU(N_i) \times U(1)$ by the condensate. The latter is an $N_i \times N_i$ complex matrix,

$$(\phi_i)_{mm} \sim \psi_{im}^1 \psi_{in}^2, \quad (8)$$

which transforms like the bi-fundamental of $SU(N_i)_1 \times SU(N_i)_2$:

$$\phi_i \rightarrow u_{i1} \phi_i u_{i2}^\dagger, \quad u_{iA} \in SU(N_i)_A. \quad (9)$$

In terms of spin-zero composites ϕ_i reads

$$\phi_i = \frac{v_i + H_i + i\Theta_i}{\sqrt{2N_i}} + (i\Pi_i^a + \Sigma_i^a) T_i^a, \quad (10)$$

where T_i^a are the $SU(N_i)$ broken generators, normalized according to $\text{Tr } T_i^a T_i^b = \delta^{ab}/2$. Here v_i is the vacuum expectation value of the condensate, and H_i is a $\psi_{im}^1 \psi_{im}^2$ scalar singlet.

If the representation R_i is real, rotations in techniflavor space *can* mix ψ^1 and ψ^2 fermions, as these transform in the same way under the TC gauge group. As a consequence the TC Lagrangian features a global $SU(2N_i)$ techniflavor symmetry, which is spontaneously broken to $SO(2N_i)$ by the condensate. The latter is a $2N_i \times 2N_i$ complex matrix,

$$(\Phi_i)_{mn}^{AB} \sim \psi_{im}^A \psi_{in}^B, \quad (11)$$

where $A, B = 1, 2$, and $m, n = 1, \dots, N_i$. The spin-zero matrix Φ_i transforms as the two-index symmetric representation of $SU(2N_i)$:

$$\Phi_i \rightarrow u_i \Phi_i u_i^T, \quad u_i \in SU(2N_i), \quad (\Phi_i)_{nm}^{BA} = (\Phi_i)_{mn}^{AB}. \quad (12)$$

In terms of spin-zero composites, Φ_i reads

$$\Phi_i = \left[\frac{v_i + H_i + i\Theta_i}{\sqrt{4N_i}} + (i\Pi_i^a + \Sigma_i^a) X_i^a \right] E_i, \quad (13)$$

where X_i^a are the broken generators belonging to the $SU(2N_i) - SO(2N_i)$ algebra, and normalized according to $\text{Tr } X_i^a X_i^b = \delta^{ab}/2$. The matrix E_i is an $SO(2N_i)$ invariant satisfying

$$E_i^T X_i^a T = X_i^a E_i, \quad E_i^T = E_i, \quad E_i E_i^\dagger = 1. \quad (14)$$

Finally, if the R_i representation is pseudoreal, the global symmetry in techniflavor space is $SU(2N_i)$, and the condensate is as in Eq. (11). However the matrix Φ_i is now in the two-index antisymmetric representation of $SU(2N_i)$,

$$\Phi_i \rightarrow u_i \Phi_i u_i^T, \quad u_i \in SU(2N_i), \quad (\Phi_i)_{nm}^{BA} = -(\Phi_i)_{mn}^{AB}, \quad (15)$$

because of an extra minus sign introduced by the invariant which contracts the TC indices (not displayed in Eq. (11)). As a consequence the condensate breaks the global symmetry to $Sp(2N_i)$ rather than $SO(2N_i)$. In terms of spin-zero composites, the $2N_i \times 2N_i$ Φ_i matrix is as in Eq. (13), where now the X_i^a broken generators belong to the $SU(2N_i) - Sp(2N_i)$ algebra, and E_i is an $Sp(2N_i)$ invariant satisfying

$$E_i^T X_i^{aT} = -X_i^a E_i, \quad E_i^T = -E_i, \quad E_i E_i^\dagger = 1. \quad (16)$$

We can unify notation for the complex and real or pseudoreal scenarios by defining, for the case of complex R_i , the $2N_i \times 2N_i$ matrix

$$\Phi_i \equiv \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & \phi_i \\ \phi_i^\dagger & 0 \end{pmatrix}. \quad (17)$$

Then, assuming no violation of techniflavor symmetry, and retaining only terms up to dimension four, the symmetry-breaking potential reads

$$\begin{aligned} V = & -\mu_1^2 \text{Tr } \Phi_1 \Phi_1^\dagger - \mu_2^2 \text{Tr } \Phi_2 \Phi_2^\dagger + \lambda'_1 \text{Tr } \Phi_1 \Phi_1^\dagger \Phi_1 \Phi_1^\dagger + \lambda''_1 \text{Tr } \Phi_1 \Phi_1^\dagger \text{Tr } \Phi_1 \Phi_1^\dagger \\ & + \lambda'_2 \text{Tr } \Phi_2 \Phi_2^\dagger \Phi_2 \Phi_2^\dagger + \lambda''_2 \text{Tr } \Phi_2 \Phi_2^\dagger \text{Tr } \Phi_2 \Phi_2^\dagger + 2\lambda \text{Tr } \Phi_1 \Phi_1^\dagger \text{Tr } \Phi_2 \Phi_2^\dagger. \end{aligned} \quad (18)$$

A few comments are in order for this potential. First, the pseudoscalars are all massless in Eq. (18). In particular, the Π_i^a fields are the Nambu-Goldstone bosons (NGBs) associated to the spontaneous breaking of the $SU(N_i)_1 \times SU(N_i)_2$ or $SU(2N_i)$ techniflavor symmetries. Three of these NGBs become the longitudinal components of the SM W and Z boson, once the electroweak interactions are “switched on”. The remaining NGBs receive mass through radiative effects and/or additional new interactions beyond TC, such as Extended TC [16, 17]. These interactions can be accounted for by adding techniflavor-breaking potential terms to Eq. (18). The Θ_i pseudoscalar singlets are massless in Eq. (18), because of additional and spontaneously broken $U(1)_i$ symmetries. These states acquire mass from instantons – in the form of $\text{Det } \Phi_i$ invariants to be added to V – and/or ETC interactions. Finally, the scalar multiplets Σ_i^a can be made heavier than the singlets (as expected from scaling up the QCD spectrum) by adjusting the quartic terms, and/or by including higher order invariant terms to the potential. We shall ignore all these issues, as our goal is to highlight the see-saw mechanism for the scalar singlets. This is fully accounted for in the potential of Eq. (18).

Minimization of the potential gives

$$v_1^2 = \frac{\mu_1^2/\lambda_1 - \lambda\mu_2^2/\lambda_1\lambda_2}{1 - \lambda^2/\lambda_1\lambda_2}, \quad v_2^2 = \frac{\mu_2^2/\lambda_2 - \lambda\mu_1^2/\lambda_1\lambda_2}{1 - \lambda^2/\lambda_1\lambda_2}, \quad (19)$$

where

$$\lambda_i \equiv \frac{\lambda'_i}{2N_i} + \lambda''_i. \quad (20)$$

The isosinglet mass Lagrangian is as in Eq. (1), with

$$M_i^2 = 2\lambda_i v_i^2, \quad (21)$$

and

$$\delta^2 = \frac{\lambda^2}{\lambda_1 \lambda_2}. \quad (22)$$

In an $SU(N_{\text{TC}})$ theory, in the limit of large N_{TC} , the double trace terms, in particular the mass mixing term, are subleading in $1/N_{\text{TC}}$, see *e.g.* [18]. To see this explicitly, consider that the contributions to λ_i and λ are dominated by the diagrams of Fig. 2 (a) and (b), respectively, where the black disks represent scalar insertions. These introduce a normalization factor $1/\sqrt{d(R_i)}$, where $d(R_i)$ is the dimension of the representation R_i . Recalling that the TC gauge coupling scales like $1/\sqrt{N_{\text{TC}}}$, we find the scaling behaviors

$$\lambda_i \sim \frac{1}{N_i d(R_i)}, \quad \lambda \sim \frac{T(R_1)T(R_2)d(G)}{d(R_1)d(R_2)N_{\text{TC}}^2}, \quad (23)$$

Here $T(R_i)$ is defined by $\text{Tr } t_i^a t_i^b = T(R_i)\delta^{ab}$, where t_i^a are the TC generators in the representation R_i , and $d(G)$ is the dimension of the TC group, $N_{\text{TC}}^2 - 1$. In the large- N_{TC} limit this gives

$$\delta^2 \sim \frac{N_1 N_2 T(R_1)^2 T(R_2)^2}{d(R_1)d(R_2)}, \quad (24)$$

For example, if R_1 is the fundamental and R_2 the adjoint representation, then $\delta^2 \sim N_1 N_2 / N_{\text{TC}}$. The fact that δ^2 decreases with N_{TC} was to be expected, as λ arises from a three-loop diagram, and is therefore subdominant in the large- N_{TC} limit. However, for small values of N_{TC} we expect δ^2 to be of order one, as suppression from loop factors is compensated by the large TC coupling. From Eq. (24) we also observe that δ^2 grows with the number of flavors.

In addition to mass mixing of scalar singlets, we also expect mass mixing of the pseudoscalar and spin-one singlets. These, however, may feature larger diagonal masses than those of the scalar singlets. We will not further address this point here.

A. Ultra-minimal technicolor

The UMT model [10] is a two-scale TC model based on an $SU(2)_{\text{TC}}$ gauge theory, with $N_1 = 2$ Dirac techniflavors in the fundamental representation, U and D , and $N_2 = 1$ Dirac techniflavor

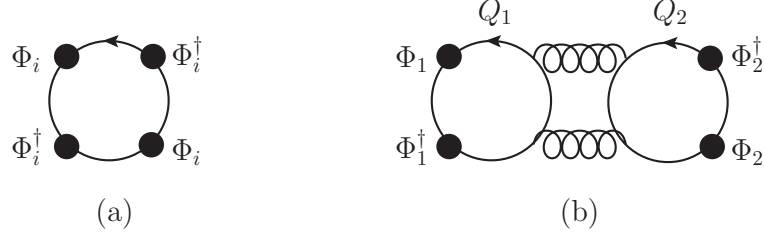


FIG. 2

in the adjoint representation, λ . The two fundamental technifermions are arranged in a doublet with respect to the weak interactions, whereas the adjoint technifermion is not charged under the electroweak interactions. UMT features the smallest contribution to the *perturbative*² electroweak S parameter, $S_{\text{pert}} \simeq 1/3\pi$, compatible with electroweak symmetry breaking and near-conformal dynamics [10, 14].

In UMT both fermion species are assumed to condense at roughly the same scale $\Lambda_1 \simeq \Lambda_2$. In fact it is readily found that $C_2(R_1) \simeq C_2(R_2)$, whence $\alpha_c(R_1) \simeq \alpha_c(R_2)$ in the ladder approximation. The technifermion-condensate vevs are $\langle \bar{U}_R U_L + \bar{D}_R D_L \rangle$ and $\langle \lambda^1 \lambda^2 \rangle$. The former breaks the electroweak symmetry and produces, among others, an isospin triplet of Goldstone bosons, $\Pi_1^{1,2,3}$. These become the longitudinal modes of the W and Z boson, which requires $v_1 = v = 246$ GeV. Based on the above assumption, we also have $v_2 \simeq v_1$. The full global symmetry breaking pattern, in the absence of electroweak interactions, is $SU(4) \times SU(2) \times U(1) \rightarrow Sp(4) \times U(1) \times Z_2$. Notice the extra $U(1)$ symmetry, relative to the general symmetry breaking patterns discussed above. This arises from the fact that a linear combination of the two $U(1)$ symmetries, in $U(4) = SU(4) \times U(1)$ and $U(2) = SU(2) \times U(1)$, is anomaly free, whereas in isolation each one of these symmetries is anomalous.

Following the above discussion we can describe the scalar sector using a linear realization of the global symmetries in terms of a 4×4 matrix Φ_1 , and a 2×2 matrix Φ_2 . Up to dimension-four terms the potential is as in Eq. (18). This gives nine massless pseudoscalars: $\Pi_1^{1,2,3,4,5}$ and $\Pi_2^{1,2}$, corresponding to $SU(4) \rightarrow Sp(4)$ and $SU(2) \rightarrow U(1)$ spontaneous symmetry breaking, respectively, plus the $\Theta_{1,2}$ pseudoscalar singlets. A linear combination of these is the Goldstone boson corresponding to the $U(1) \rightarrow Z_2$ spontaneous symmetry breaking, whereas the remaining linear combination receives mass from instantons, in the form of higher dimensional terms. The one of lowest order has

² By S_{pert} we mean the computation of S from a loop of technifermions Q with a dynamical mass $m_Q \gg m_Z$, see *e.g.* the discussion in [19].

dimension six:

$$(\det \Phi_2)^2 \text{Pf } \Phi_1 + \text{h.c.} . \quad (25)$$

Incidentally this provides an additional mass mixing term for the two scalar singlets. The UMT model is thus a prime TC example where the lightest scalar mass eigenstate can be as light as 125 GeV, as exemplified by the black curve in Fig. 1.

Note however that before mass mixing the scalar masses in the UMT model might be well below the TeV scale due to the argued walking dynamics of the model. For example the mass of the lightest scalar in the UMT model has been estimated to be as low as 250 GeV in Ref. [20]. This model computation does not take into account mass mixing between the scalars, but relies on the assumed near-conformality of the theory.

B. Low-scale TC

In the two-scale TC framework proposed early on in Ref. [8] the dynamical assumption is that $\Lambda_1 \ll \Lambda_2$ and such models are also referred to as low-scale TC [9]. This hierarchy in scales requires choosing the representations such that the quadratic Casimirs satisfy $C_2(R_1) \ll C_2(R_2)$ and/or lead to a small $\beta_0(R_1)$ ³. The low-scale TC assumption that both sectors have the technifermions arranged in weak doublets implies that the electroweak scale v must be related to v_i via $v = \sqrt{N_1 v_1^2 + N_2 v_2^2}$ to ensure the correct W and Z masses. For non-large values of N_1 , $v \simeq \sqrt{N_2} v_2$.

Again we can describe the scalar sector using a linear realization of the global symmetries in terms of appropriate matrices of composite fields Φ_i . The diagonal mass M_1 is by construction relatively light [21], and mass mixing will further reduce its value. The corresponding scenario is similar to the one depicted by the red curves in Fig. 1.

IV. SUMMARY AND DISCUSSION

In this paper we have discussed how a light composite scalar may arise in two-scale TC theories via mass mixing between relatively heavy scalar resonances. We have argued that the mass of the light scalar may be compatible with the recently observed ~ 125 GeV resonance and discussed a concrete minimal two-scale TC model, UMT [14], which provides mass mixing of the required order of magnitude. We have also discussed the mechanism in low-scale TC models [8], where

³ Achieving this typically requires a large number of technifermions in the representation R_1 , such as the fundamental, or alternatively four-technifermion operators with large enough coefficients [15].

the lightest scalar resonance is expected to be relatively light compared to the TeV scale, already before including effects of mass-mixing. Finally radiative corrections from the interaction with the top quark can further reduce the mass of the lightest scalar resonance [13].

It will be interesting to study the phenomenology of TC models featuring a scalar see-saw mechanism in the light of LHC data, already indicating that the couplings of the scalar resonance to the SM fermions and gauge bosons must be SM Higgs-like.

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